

Baryonic Rare Decays of $\Lambda_b \rightarrow \Lambda \ell^+ \ell^-$

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Abstract

We present a systematic analysis for the rare baryonic exclusive decays of $\Lambda_b \rightarrow \Lambda \ell^+ \ell^-$ ($\ell = e, \mu, \tau$). We study the differential decay rates and the di-lepton forward-backward, lepton polarization and various CP asymmetries with a new simple set of form factors inspired by the heavy quark effective theory. We show that most of the observables are insensitive to the non-perturbative QCD effects. To illustrate the effect of new physics, we discuss our results in an explicit supersymmetric extension of the standard model, which contains new CP violating phases and therefore induces sizable CP violating asymmetries.

1 Introduction

A priority in current particle physics research is to determine the parameters of the Cabibbo-Kobayashi-Maskawa (CKM) matrix elements [1] in the standard model (SM). Due to the CLEO measurement of the radiative $b \rightarrow s\gamma$ decay [2], some interest has been focused on the rare decays related to $b \rightarrow sl^+l^-$ induced by the flavor changing neutral currents (FCNCs). In the SM, these rare decays occur at the loop level and depend on the CKM elements. In the literature, most of studies have been concentrated on the corresponding exclusive rare B-meson decays such as $B \rightarrow K^{(*)}l^+l^-$ [3]. However, these exclusive modes contain several unknown hadronic form factors, which cannot be measured in the present B-meson facilities unlike the kaon cases. Recently, we have examined the exclusive rare baryonic decays of $\Lambda_b \rightarrow \Lambda l\bar{l}$ ($l = \nu, e, \mu, \tau$) [4, 5, 6] and found that some of physical quantities are insensitive to the hadronic uncertainties.

In this paper, we give a systematic study on the baryonic decays of $\Lambda_b \rightarrow \Lambda l^+l^-$. We will explore various possible CP even and odd asymmetries to show how the hadronic unknown parameters are factored out in most of cases. To illustrate CP violating effect, we will also discuss an explicit CP violating model with SUSY.

The paper is organized as follows. In Sec. 2, we give the effective Hamiltonian for the decays of $\Lambda_b \rightarrow \Lambda l\bar{l}$ and the most general form factors in the $\Lambda_b \rightarrow \Lambda$ transition. In Sec. 3, we derive the general forms of the differential decay rates. In Sec. 4, we study the di-lepton forward-backward, lepton polarization and various CP violating asymmetries. We perform our numerical analysis in Sec. 5. We present our conclusions in Sec. 6.

2 Effective Hamiltonian and form factors

In the SM, the effective Hamiltonian for $b \rightarrow sl^+l^-$ is given by

$$\mathcal{H} = -4 \frac{G_F}{\sqrt{2}} V_{tb} V_{ts}^* \sum_{i=1}^{10} C_i(\mu) O_i(\mu) \quad (1)$$

where the expressions of the renormalized Wilson coefficients $C_i(\mu)$ and operators $O_i(\mu)$ can be found in Ref. [7]. From Eq. (1), the free quark decay amplitude is written as

$$\begin{aligned} \mathcal{M}(b \rightarrow sl^+l^-) &= \frac{G_F \alpha_{em}}{\sqrt{2}\pi} V_{tb} V_{ts}^* \left[\bar{s} \left(C_9^{eff}(\mu) \gamma_\mu P_L - \frac{2m_b}{q^2} C_7^{eff}(\mu) i\sigma_{\mu\nu} q^\nu P_R \right) b \bar{l} \gamma^\mu l \right. \\ &\quad \left. + \bar{s} C_{10} \gamma_\mu P_L b \bar{l} \gamma^\mu \gamma_5 l \right] \end{aligned} \quad (2)$$

with $P_{L(R)} = (1 \mp \gamma_5)/2$. We note that in Eq. (2), only the term associated with the Wilson coefficient C_{10} is independent of the μ scale. Besides the short-distance (SD) contributions, the long-distance (LD) ones such as that from the $c\bar{c}$ resonant states of $\Psi, \Psi' \dots etc$ are also important for the decay rate. It is known that for the LD effects in the B-meson decays [8, 9, 10, 11, 12, 13], both the factorization assumption (FA) and the vector meson dominance (VMD) approximation have been used. In baryonic decays, we assume that the parametrization of LD contributions is the same as that in the B-meson decays. Hence, we may include the resonant effect (RE) by absorbing it to the corresponding Wilson coefficient. In this paper as a more complete analysis we also include the LD contributions to the decay of $b \rightarrow s\gamma$, induced by the nonfactorizable

effects [14, 15]. The effective Wilson coefficients of C_9^{eff} and C_7^{eff} can be expressed as the standard form

$$C_9^{eff}(\mu) = C_9(\mu) + Y(z, s'), \quad (3)$$

$$C_7^{eff}(\mu) = C_7(\mu) + C_7'(\mu, q^2), \quad (4)$$

where

$$\begin{aligned} Y(z, s') &= \left(h(z, s') + \frac{3}{\alpha_{em}^2} \sum_{j=\Psi, \Psi'} k_j \frac{\pi \Gamma(j \rightarrow l^+ l^-) M_j}{q^2 - M_j^2 + i M_j \Gamma_j} \right) \\ &\quad - \frac{1}{2} h(1, s') (4C_3 + 4C_4 + 3C_5 + C_6) - \frac{1}{2} h(0, s') (C_3 + 3C_4), \\ C_7'(\mu, q^2) &= C_{b \rightarrow s\gamma}'(\mu) + \omega \left(h(z, s') + \frac{3}{\alpha_{em}^2} \sum_{j=\Psi, \Psi'} k_j \frac{\pi \Gamma(j \rightarrow l^+ l^-) M_j}{q^2 - M_j^2 + i M_j \Gamma_j} \right), \end{aligned} \quad (5)$$

with

$$\begin{aligned} h(z, s') &= -\frac{8}{9} \ln z + \frac{8}{27} + \frac{4}{9}x - \frac{2}{9}(2+x)|1-x|^{1/2} \\ &\quad \times \begin{cases} \ln \left| \frac{\sqrt{1-x}+1}{\sqrt{1-x}-1} \right| - i\pi & \text{for } x \equiv 4z^2/s' < 1 \\ 2 \arctan \frac{1}{\sqrt{x-1}} & \text{for } x \equiv 4z^2/s' > 1 \end{cases}, \\ C_{b \rightarrow s\gamma}' &= i\alpha_s \left[\frac{2}{9} \eta^{14/23} (G_1(x_t) - 0.1687) - 0.03 C_2(\mu) \right], \\ G_1(x) &= \frac{x(x^2 - 5x - 2)}{8(x-1)^3} + \frac{3x^2 \ln x}{4(x-1)^4}. \end{aligned} \quad (6)$$

Here $Y(z, s')$ combines the one-loop matrix elements and the LD contributions of operators $O_1 - O_6$, $C_{b \rightarrow s\gamma}'$ is the absorptive part of $b \rightarrow s\gamma$ [16] with neglecting the small contribution from $V_{ub}V_{us}^*$, $z = m_c/m_b$, $s' = q^2/m_b^2$, $\eta = \alpha_s(m_W)/\alpha_s(\mu)$, $x_t = m_t^2/m_W^2$, M_j (Γ_j) are the masses (widths) of intermediate states, $|\omega| \leq 0.15$ describing the nonfactorizable contributions to $b \rightarrow s\gamma$ decay at $q^2 = 0$ [14, 15], and the factors k_j are phenomenological parameters for compensating the approximations of the FA and VMD and reproducing the correct branching ratios of $B(\Lambda_b \rightarrow \Lambda J/\Psi \rightarrow \Lambda l^+ l^-) = B(\Lambda_b \rightarrow \Lambda J/\Psi) \times B(J/\Psi \rightarrow l^+ l^-)$ when we study the Λ_b decays. We note that by taking $k_\Psi \simeq -1/(3C_1 + C_2)$ and $B(\Lambda_b \rightarrow \Lambda J/\Psi) = (4.7 \pm 2.8) \times 10^{-4}$, the k_j factors in the Λ_b case are almost the same as that in the B-meson one [5]. The Wilson coefficients (WCs) at the scale of $\mu \sim m_b \sim 4.8$ GeV are shown in Table 1.

Table 1: Wilson coefficients for $m_t = 170$ GeV, $\mu = 4.8$ GeV.

WC	C_1	C_2	C_3	C_4	C_5
	-0.226	1.096	0.01	-0.024	0.007
WC	C_6	C_7	C_8	C_9	C_{10}
	-0.028	-0.305	-0.15	4.186	-4.559

Using the form factors given in Appendix A, we write the amplitude of $\Lambda_b \rightarrow \Lambda l^+ l^-$ as

$$\mathcal{M}(\Lambda_b \rightarrow \Lambda l^+ l^-) = \frac{G_F \alpha_{em}}{\sqrt{2}\pi} V_{tb} V_{ts}^* \{H_1^\mu L_\mu + H_2^\mu L_\mu^5\} \quad (7)$$

where

$$\begin{aligned} H_1^\mu &= \bar{\Lambda} \gamma^\mu (A_1 P_R + B_1 P_L) \Lambda_b + \bar{\Lambda} i \sigma^{\mu\nu} q_\nu (A_2 P_R + B_2 P_L) \Lambda_b, \\ H_2^\mu &= \bar{\Lambda} \gamma^\mu (D_1 P_R + E_1 P_L) \Lambda_b + \bar{\Lambda} i \sigma^{\mu\nu} q_\nu (D_2 P_R + E_2 P_L) \Lambda_b \\ &\quad + q^\mu \bar{\Lambda} (D_3 P_R + E_3 P_L) \Lambda_b, \end{aligned} \quad (8)$$

$$\begin{aligned} L_\mu &= \bar{l} \gamma_\mu l, \\ L_\mu^5 &= \bar{l} \gamma_\mu \gamma_5 l \end{aligned} \quad (9)$$

with

$$\begin{aligned} A_i &= C_9^{eff} \frac{f_i - g_i}{2} - \frac{2m_b}{q^2} C_7^{eff} \frac{f_i^T + g_i^T}{2}, \\ B_i &= C_9^{eff} \frac{f_i + g_i}{2} - \frac{2m_b}{q^2} C_7^{eff} \frac{f_i^T - g_i^T}{2}, \\ D_i &= C_{10} \frac{f_i - g_i}{2}, \\ E_i &= C_{10} \frac{f_i + g_i}{2}. \end{aligned} \quad (10)$$

and $i = 1, 2, 3$.

The processes for the heavy to light baryonic decays such as those with $\Lambda_b \rightarrow \Lambda$ have been studied based on the heavy quark effective theory (HQET) in Ref. [17] and it is found that

$$\langle \Lambda(p_\Lambda) | \bar{s} \Gamma b | \Lambda_b(p_{\Lambda_b}) \rangle = \bar{u}_\Lambda \left(F_1(q^2) + \not{q} F_2(q^2) \right) \Gamma u_{\Lambda_b} \quad (11)$$

where Γ denotes the Dirac matrix, $v = p_{\Lambda_b}/M_{\Lambda_b}$ is the four-velocity of Λ_b , $q = p_{\Lambda_b} - p_\Lambda$ is the momentum transfer, and $F_{1,2}$ are the form factors. Clearly, there are only two independent form factors $F_{1,2}$ in the HQET. Comparing with the general forms of the form factors in Appendix A, we get the relations among the form factors as follows:

$$\begin{aligned} g_1 &= f_1 = f_2^T = g_2^T = F_1 + \sqrt{r} F_2, \\ g_2 &= f_2 = g_3 = f_3 = g_T^V = f_T^V = \frac{F_2}{M_{\Lambda_b}}, \\ g_T^S &= f_T^S = 0, \\ g_1^T &= f_1^T = \frac{F_2}{M_{\Lambda_b}} q^2, \\ g_3^T &= \frac{F_2}{M_{\Lambda_b}} (M_{\Lambda_b} + M_\Lambda), \quad f_3^T = -\frac{F_2}{M_{\Lambda_b}} (M_{\Lambda_b} - M_\Lambda), \end{aligned} \quad (12)$$

where $r = M_\Lambda^2/M_{\Lambda_b}^2$. From the CLEO result of $R = -0.25 \pm 0.14 \pm 0.08$ [18], we know that $|F_2| < |F_1|$. Due to Eq. (12), only f_1 (g_1) and f_2^T (g_2^T) are proportional to F_1 and therefore,

they are large, whereas all the others are small since they are related to the small form factor F_2 . Furthermore, from Eq. (10), we find that $\{f^T\}$ and $\{g^T\}$ are associated with C_7 which is about one order of the magnitude smaller than C_9 and C_{10} so that their effects to the deviation of the results in the HQET are small. Hence, with the information of the HQET, we can make a good approximation for the general form factors of transition matrix elements given in Eqs. (7) and (10). Altogether, we have the following relations:

$$\begin{aligned}\bar{f} &\equiv \frac{f_1 + g_1}{2}, \quad \frac{f_2^T + g_2^T}{f_1 + g_1} \simeq 1 \\ \frac{f_1 - g_1}{f_1 + g_1} &\simeq \delta, \quad \frac{g_2}{f_2} \simeq \frac{g_1^T}{f_1^T} \simeq \frac{g_2^T}{f_2^T} \simeq 1, \\ \frac{f_1^T + g_1^T}{f_1 + g_1} \frac{1}{q^2} &\simeq \frac{f_2 + g_2}{f_1 + g_1}.\end{aligned}\tag{13}$$

In the HQET, it is easy to show that

$$\delta = 0, \quad \rho \equiv M_{\Lambda_b} \left(\frac{f_2 + g_2}{f_1 + g_1} \right) = \frac{F_2}{F_1 + \sqrt{r}F_2}.\tag{14}$$

3 Differential decays rates

In this section we first present the formulas by including the lepton mass for the double differential decay rates with respect to the angle of the lepton and the invariant mass of the di-lepton. In the following we only show the results of the SM with the form factors in Eq. (13). The general ones with including right-handed couplings are presented in Appendix B.

Introducing dimensionless variables of $t = p_{\Lambda_b} \cdot p_{\Lambda} / M_{\Lambda_b}^2$, $r = M_{\Lambda}^2 / M_{\Lambda_b}^2$, $\hat{m}_l = m_l / M_{\Lambda_b}$, $\hat{m}_b = m_b / M_{\Lambda_b}$, and $s = q^2 / M_{\Lambda_b}^2$, the double partial differential decay rates for $\Lambda_b \rightarrow \Lambda l^+ l^-$ ($l = e, \mu, \tau$) can be written as

$$\frac{d^2\Gamma}{dsd\hat{z}} = \frac{G_F^2 \alpha_{em}^2 \lambda_t^2}{768\pi^5} M_{\Lambda_b}^5 \sqrt{\phi(s)} \sqrt{1 - \frac{4m_l^2}{q^2}} \bar{f}^2 R_{\Lambda_b}(s, \hat{z}),\tag{15}$$

where

$$R_{\Lambda_b}(s, \hat{z}) = I_0(s, \hat{z}) + \hat{z} I_1(s, \hat{z}) + \hat{z}^2 I_2(s, \hat{z})\tag{16}$$

and

$$\begin{aligned}I_0(s, \hat{z}) &= -6\sqrt{r}s \left[-2\hat{m}_b\rho \left(1 + 2\frac{m_l^2}{q^2} \right) \text{Re}C_9^{\text{eff}} C_7^{\text{eff}*} \right. \\ &\quad \left. + \delta \left(\left(1 + 2\frac{m_l^2}{q^2} \right) |C_9^{\text{eff}}|^2 + \left(1 - 6\frac{m_l^2}{q^2} \right) |C_{10}|^2 \right) \right] \\ &\quad + \frac{3}{4} \left((1-r)^2 - s^2 \right) \left[(2\hat{m}_b\rho)^2 |C_7^{\text{eff}}|^2 + |C_9^{\text{eff}}|^2 + |C_{10}|^2 \right] \\ &\quad + 6\hat{m}_l^2 t \left[(2\hat{m}_b\rho)^2 |C_7^{\text{eff}}|^2 + |C_9^{\text{eff}}|^2 - |C_{10}|^2 \right]\end{aligned}$$

$$\begin{aligned}
& +6\sqrt{r}(1-t) \left\{ 4 \left(1 + 2\frac{m_l^2}{q^2} \right) \hat{m}_b^2 \rho |C_7^{eff}|^2 \right. \\
& \left. + \rho s \left[\left(1 + 2\frac{m_l^2}{q^2} \right) |C_9^{eff}|^2 + \left(1 - 2\frac{m_l^2}{q^2} \right) |C_{10}|^2 \right] \right\} \\
& +12 \left(1 + 2\frac{m_l^2}{q^2} \right) \hat{m}_b (t-r) (1+s\rho^2) \text{Re}C_9^{\text{eff}} C_7^{\text{eff}*} \\
& +12 \left(1 + 2\frac{m_l^2}{q^2} \right) \hat{m}_b \sqrt{r} s \rho \text{Re}C_9^{\text{eff}} C_7^{\text{eff}*} \\
& -6 \left[s(1-t)(t-r) - \frac{1}{8} ((1-r)^2 - s^2) \right] \\
& \times \left[\frac{4\hat{m}_b^2}{s} |C_7^{eff}|^2 + s\rho^2 (|C_9^{eff}|^2 + |C_{10}|^2) \right] \\
& -6\hat{m}_l^2 (2r - (1+r)t) \left[\left(\frac{2\hat{m}_b}{s} \right)^2 |C_7^{eff}|^2 + \rho^2 (|C_9^{eff}|^2 - |C_{10}|^2) \right], \quad (17)
\end{aligned}$$

$$\begin{aligned}
I_1(s, \hat{z}) &= 3\sqrt{1 - \frac{4m_l^2}{q^2}} \phi(s) \left\{ s \left[1 - 2\sqrt{r}\rho - (1-r)\rho^2 \right] \text{Re}C_9^{\text{eff}} C_{10}^* \right. \\
&\quad \left. + 2\hat{m}_b (1 - s\rho^2) \text{Re}C_7^{\text{eff}} C_{10}^* \right\}, \quad (18)
\end{aligned}$$

$$\begin{aligned}
I_2(s, \hat{z}) &= -\frac{3}{4} \phi(s) \left(1 - 4\frac{m_l^2}{q^2} \right) \left[(2\hat{m}_b \rho)^2 |C_7^{eff}|^2 + |C_9^{eff}|^2 + |C_{10}|^2 \right] \\
&\quad + \frac{3}{4} \phi(s) \left(1 - 4\frac{m_l^2}{q^2} \right) \left[\frac{4\hat{m}_b^2}{s} |C_7^{eff}|^2 + s\rho^2 (|C_9^{eff}|^2 + |C_{10}|^2) \right], \quad (19)
\end{aligned}$$

with $\hat{z} = \hat{p}_B \cdot \hat{p}_{l^+}$ being the angle between the momenta of Λ_b and l^+ in the di-lepton invariant mass frame and $\phi(s) = (1-r)^2 - 2s(1+r) + s^2$. Here, for the simplicity, we have not displayed the dependence of the μ scale in effective Wilson coefficients. We note that the main nonperturbative QCD effect from \bar{f} has been factored out in Eq. (15). The function $R_{\Lambda_b}(s, \hat{z})$ is only related to the two parameters of δ and ρ which become one in the HQET. Since ρ is the ratio of form factors and insensitive to the QCD models, the QCD effects in the baryonic di-lepton decays are clearly less significant. Therefore, these decay modes are good physical observable to test the SM.

After integrating the angular dependence, the invariant mass distributions as function of s are given by

$$\frac{d\Gamma(\Lambda_b \rightarrow \Lambda l^+ l^-)}{ds} = \frac{G_F^2 \alpha_{em}^2 \lambda_t^2}{384\pi^5} M_{\Lambda_b}^5 \sqrt{\phi(s)} \sqrt{1 - \frac{4m_l^2}{q^2}} \bar{f}^2 R_{\Lambda_b}(s), \quad (20)$$

where

$$R_{\Lambda_b}(s) = \Gamma_1(s) + \Gamma_2(s) + \Gamma_3(s) \quad (21)$$

with

$$\Gamma_1(s) = -6\sqrt{r}s \left[-2\hat{m}_b \rho \left(1 + 2\frac{m_l^2}{q^2} \right) \text{Re}C_9^{\text{eff}} C_7^{\text{eff}*} \right]$$

$$\begin{aligned}
& +\delta \left(\left(1 + 2\frac{m_l^2}{q^2} \right) |C_9^{eff}|^2 + \left(1 - 6\frac{m_l^2}{q^2} \right) |C_{10}|^2 \right) \Big] \\
& + \left[-2r \left(1 + 2\frac{m_l^2}{q^2} \right) - 4t^2 \left(1 - \frac{m_l^2}{q^2} \right) + 3(1+r)t \right] \\
& \times \left[(2\hat{m}_b\rho)^2 |C_7^{eff}|^2 + |C_9^{eff}|^2 + |C_{10}|^2 \right] \\
& + 6\hat{m}_l^2 t \left[(2\hat{m}_b\rho)^2 |C_7^{eff}|^2 + |C_9^{eff}|^2 - |C_{10}|^2 \right], \tag{22}
\end{aligned}$$

$$\begin{aligned}
\Gamma_2(s) = & 6\sqrt{r}(1-t) \left\{ 4 \left(1 + 2\frac{m_l^2}{q^2} \right) \hat{m}_b^2 \rho |C_7^{eff}|^2 \right. \\
& + \rho s \left[\left(1 + 2\frac{m_l^2}{q^2} \right) |C_9^{eff}|^2 + \left(1 - 2\frac{m_l^2}{q^2} \right) |C_{10}|^2 \right] \Big\} \\
& + 12 \left(1 + 2\frac{m_l^2}{q^2} \right) \hat{m}_b (t-r) (1+s\rho^2) \text{Re} C_9^{\text{eff}} C_7^{\text{eff}*}, \tag{23}
\end{aligned}$$

$$\begin{aligned}
\Gamma_3(s) = & 12 \left(1 + 2\frac{m_l^2}{q^2} \right) \hat{m}_b \sqrt{r} s \rho \text{Re} C_9^{\text{eff}} C_7^{\text{eff}*} \\
& - \left[2t^2 \left(1 + 2\frac{m_l^2}{q^2} \right) + 4r \left(1 - \frac{m_l^2}{q^2} \right) - 3(1+r)t \right] \\
& \times \left[\frac{4\hat{m}_b^2}{s} |C_7|^2 + s\rho^2 \left(|C_9^{eff}|^2 + |C_{10}|^2 \right) \right] \\
& - 6\hat{m}_l^2 (2r - (1+r)t) \left[\left(\frac{2\hat{m}_b}{s} \right)^2 |C_7^{\text{eff}}|^2 + \rho^2 \left(|C_9^{\text{eff}}|^2 - |C_{10}|^2 \right) \right]. \tag{24}
\end{aligned}$$

The limits for s are given by

$$4\hat{m}_l^2 \leq s \leq (1 - \sqrt{r})^2. \tag{25}$$

From Eqs. (22)-(24), we see that ρ appears either as $\sqrt{r}\rho$ or ρ^2 which is small since $r \sim 0.04$ and $|\rho| \sim 0.25$.

4 Lepton and CP asymmetries

4.1 Forward-backward asymmetries

The differential and normalized forward-backward asymmetries (FBAs) for the decays of $\Lambda_b \rightarrow \Lambda l^+ l^-$ as a function of s are defined by

$$\frac{dA_{FB}(s)}{ds} = \left[\int_0^1 d\hat{z} \frac{d^2\Gamma(s, \hat{z})}{ds d\hat{z}} - \int_{-1}^0 d\hat{z} \frac{d^2\Gamma(s, \hat{z})}{ds d\hat{z}} \right] \tag{26}$$

and

$$\mathcal{A}_{FB}(s) = \frac{1}{d\Gamma(s)/ds} \left[\int_0^1 d\hat{z} \frac{d^2\Gamma(s, \hat{z})}{ds d\hat{z}} - \int_{-1}^0 d\hat{z} \frac{d^2\Gamma(s, \hat{z})}{ds d\hat{z}} \right], \tag{27}$$

respectively. Explicitly, using Eq. (15), we obtain

$$\frac{dA_{FB}(s)}{ds} = \frac{G_F^2 \alpha_{em}^2 \lambda_t^2}{2^8 \pi^5} M_{\Lambda_b}^5 \phi(s) \left(1 - 4 \frac{\hat{m}_l^2}{\hat{s}}\right) \bar{f}^2 R_{FB}(s) \quad (28)$$

and

$$\mathcal{A}_{FB}(s) = \frac{3}{2} \sqrt{\phi(s)} \sqrt{1 - \frac{4\hat{m}_l^2}{s} \frac{R_{FB}(s)}{R_{\Lambda_b}(s)}} \quad (29)$$

where

$$\begin{aligned} R_{FB}(s) = & s \left[1 - 2\sqrt{r}\rho - (1-r)\rho^2\right] \text{Re}C_9^{\text{eff}}C_{10}^* \\ & + 2\hat{m}_b(1-s\rho^2) \text{Re}C_7^{\text{eff}}C_{10}^*. \end{aligned} \quad (30)$$

It is known that the FBA is a parity-odd but CP-even observable, which depends on the chirality of the leptonic and hadronic currents. In order to obtain one power of \hat{z} dependence, the related differential decay rate should be associated with $\text{Tr}L_\mu L_\nu^5$. This explains why the FBAs depend on $\text{Re}C_9^{\text{eff}}C_{10}^*$ and $\text{Re}C_7^{\text{eff}}C_{10}^*$. However, unlike that in the decays of $B \rightarrow Kl^+l^-$ where the FBAs are always zero since they only involve vector and tensor types of currents, the transition matrix elements in the baryonic decays preserve the chirality of free quark interaction.

Similar to the B-meson decays [3, 19] the FBA in Eq. (29) vanishes at s_0 which satisfies with the relation

$$\text{Re}C_9^{\text{eff}}C_{10}^* = -\frac{2\hat{m}_b}{s_0} \frac{1-s_0\rho^2}{1-2\sqrt{r}\rho-(1-r)\rho^2} \text{Re}C_7^{\text{eff}}C_{10}^*. \quad (31)$$

We will see later that the vanishing point is only sensitive to the effects of weak interaction.

4.2 Lepton polarization asymmetries

To display the spin effects of the lepton, we choose the four-spin vector of l^+ in terms of a unit vector, $\hat{\xi}$, along the spin of l^+ in its rest frame, as

$$s_+^0 = \frac{\vec{p}_+ \cdot \hat{\xi}}{m_l}, \quad \vec{s}_+ = \hat{\xi} + \frac{s_+^0}{E_{l^+} + m_l} \vec{p}_+, \quad (32)$$

and the unit vectors along the longitudinal and transverse components of the l^+ polarization to be

$$\begin{aligned} \hat{e}_L &= \frac{\vec{p}_+}{|\vec{p}_+|}, \\ \hat{e}_T &= \frac{\vec{p}_\Lambda \times \vec{p}_+}{|\vec{p}_\Lambda \times \vec{p}_+|}, \\ \hat{e}_N &= \hat{e}_L \times \hat{e}_T, \end{aligned} \quad (33)$$

respectively.

Defining the longitudinal and transverse l^+ polarization asymmetries by

$$P_i(\hat{s}) = \frac{d\Gamma(\hat{e}_i \cdot \hat{\xi} = 1) - d\Gamma(\hat{e}_i \cdot \hat{\xi} = -1)}{d\Gamma(\hat{e}_i \cdot \hat{\xi} = 1) + d\Gamma(\hat{e}_i \cdot \hat{\xi} = -1)}, \quad (34)$$

with $i = L$ and T , we find that

$$P_L = \sqrt{1 - \frac{4m_l^2}{q^2} \frac{R_L(s)}{R_{\Lambda_b}(s)}}, \quad (35)$$

$$P_T = \frac{3}{4}\pi\hat{m}_l \sqrt{1 - \frac{4m_l^2}{q^2}} \sqrt{s\phi(s)} \frac{R_T(s)}{R_{\Lambda_b}(s)}, \quad (36)$$

where

$$\begin{aligned} R_L = & -\text{Re}C_9^{\text{eff}}C_{10}^* \left[(1-r)^2 + s(1+r) - 2s^2 + 6\sqrt{r}\rho s(1-r+s) \right. \\ & \left. + \rho^2 s \left(2(1-r)^2 - s(1+r) - s^2 \right) \right] \\ & - 6\hat{m}_b \text{Re}C_7^{\text{eff}}C_{10}^* \left[(1-r-s) \left(1 + \rho^2 s \right) + 4\sqrt{r}\rho s \right], \end{aligned} \quad (37)$$

$$\begin{aligned} R_T = & \left[1 - 2\sqrt{r}\rho - \rho^2(1-r) \right] \text{Im}C_9^{\text{eff}}C_{10}^* \\ & + \frac{2\hat{m}_b}{s} \left(1 - \rho^2 s \right) \text{Im}C_7^{\text{eff}}C_{10}^*. \end{aligned} \quad (38)$$

Here we do not discuss the normal polarization (P_N) because the nonperturbative effects from the form factors are large at the small s region and moreover, the dependence of Wilson coefficients is similar to the invariant mass distribution [6]. We note that the longitudinal lepton polarization of P_L in Eq. (35) is also a parity-odd and CP-even observable just like the FBA, whereas P_T in Eq. (36) a T-odd one which is related to the triple correlation of $\vec{s}_+ \cdot (\vec{p}_\Lambda \times \vec{p}_+)$. In general, P_T can be induced without CP violation as the cases in B-meson [20] and kaon [21] decays. However, we expect that they are small. Moreover, such effects can be extracted away while we consider the difference between the particle and anti-particle as discussed in the next section.

4.3 CP asymmetries

In this subsection, we define the following interesting direct CP asymmetries (CPAs) by

$$\Delta_\Gamma = \frac{d\Gamma - d\bar{\Gamma}}{d\Gamma + d\bar{\Gamma}}, \quad (39)$$

$$\Delta_{FB} = \frac{d\Gamma_{FB} - d\bar{\Gamma}_{FB}}{d\Gamma + d\bar{\Gamma}}, \quad (40)$$

$$\Delta_{P_i} = \frac{d\Gamma(\vec{\xi} \cdot \vec{e}_i) - d\bar{\Gamma}(\vec{\xi} \cdot \vec{e}_i)}{d\Gamma + d\bar{\Gamma}}, \quad i = L, T \quad (41)$$

where we have used $d\Gamma + d\bar{\Gamma}$ as the normalization. The above four CPAs are CP-odd quantities and they are CP violating observables. For Δ_{Γ, FB, P_L} in Eqs. (39)-(41), to

display the difference of the physical observable between the particle and anti-particle, it is necessary to have the strong and weak phases simultaneously in the processes. In the decays of $b \rightarrow sl^+l^-$ ($l = e, \mu, \tau$), the strong phases are generated by the absorptive parts of one-loop matrix elements in operators $O_1 \sim O_6$ and LD contributions. However, since P_T is a T -odd observable and only related to the imaginary couplings, even without strong phases, we still can have nonzero values of $\Delta_T(\Lambda_b \rightarrow \Lambda l^+l^-)$. For $b \rightarrow sl^+l^-$ and $\bar{b} \rightarrow \bar{s}l^-l^+$ decays the Wilson coefficients $C_9^{eff}(\mu)$ and $C_7^{eff}(\mu)$ in Eqs. (3) and (4) can be rewritten as

$$\begin{aligned} C_9^{eff}(\mu) &= C_9^0(\mu) + iC_9^{abs}(\mu), \\ \bar{C}_9^{eff}(\mu) &= C_9^{0*}(\mu) + iC_9^{abs}(\mu), \\ C_7^{eff}(\mu) &= C_7^0(\mu) + iC_7^{abs}(\mu), \\ \bar{C}_7^{eff}(\mu) &= C_7^{0*}(\mu) + iC_7^{abs}(\mu), \end{aligned} \quad (42)$$

with

$$\begin{aligned} C_9^0(\mu) &= C_9(\mu) + \text{Re}Y(z, s'), \\ C_7^0(\mu) &= C_7(\mu) + \text{Re}C_7'(\mu, q^2), \\ C_9^{abs}(\mu) &= \text{Im}Y(z, s'), \\ C_7^{abs}(\mu) &= \text{Im}C_7'(\mu, q^2), \end{aligned} \quad (43)$$

where we have assumed that the strong phases are all from the SM and there are no weak phases in absorptive parts. We note that there is no strong phase in C_{10} .

According to Eqs. (20), (29), (35), and (36), the CP asymmetries are all related to the following combinations:

$$\begin{aligned} \text{Re}C_9^{eff}C_7^{eff*} - \text{Re}\bar{C}_9^{eff}\bar{C}_7^{eff*} &= 2C_9^{abs}\text{Im}C_7^0 + 2C_7^{abs}\text{Im}C_9^0, \\ \text{Re}C_{7,9}^{eff}C_{10}^* - \text{Re}\bar{C}_{7,9}^{eff}\bar{C}_{10}^* &= 2C_{7,9}^{abs}\text{Im}C_{10}, \\ \text{Im}C_{7,9}^{eff}C_{10}^* - \text{Im}\bar{C}_{7,9}^{eff}\bar{C}_{10}^* &= 2\text{Im}C_{7,9}^0C_{10}^* + 2C_{7,9}^{abs}\text{Im}C_{10}, \\ |C_{7,9}^{eff}|^2 - |\bar{C}_{7,9}^{eff}|^2 &= 4C_{7,9}^{abs}\text{Im}C_{7,9}^0. \end{aligned} \quad (44)$$

Explicitly, the CP asymmetries in Eqs. (39), (40), and (41) are found to be

$$\begin{aligned} \Delta_\Gamma &= \frac{2}{R_{\Lambda_b}(s)} \left\{ 6\hat{m}_b \left(1 + 2\frac{m_l^2}{q^2} \right) \left[2\sqrt{r}\rho s + (t-r)(1+s\rho^2) \right] \left[C_9^{abs}\text{Im}C_7^0 + C_7^{abs}\text{Im}C_9^0 \right] \right. \\ &\quad + \left[-2r \left(1 + 2\frac{m_l^2}{q^2} \right) - 4t^2 \left(1 - \frac{m_l^2}{q^2} \right) + 3(1+r)t + 6\hat{m}_l^2 t \right] \\ &\quad \times \left[4\hat{m}_b^2 \rho^2 C_7^{abs}\text{Im}C_7^0 + C_9^{abs}\text{Im}C_9^0 \right] \\ &\quad + \left[-2t^2 \left(1 + 2\frac{m_l^2}{q^2} \right) - 4r \left(1 - \frac{m_l^2}{q^2} \right) + 3(1+r)t - 6\frac{\hat{m}_l^2}{s} (2r - t - tr) \right] \\ &\quad \times \left[\frac{4\hat{m}_b^2}{s} C_7^{abs}\text{Im}C_7^0 + s\rho^2 C_9^{abs}\text{Im}C_9^0 \right] + 6\sqrt{r}\rho(1-t) \left(1 + 2\frac{m_l^2}{q^2} \right) \\ &\quad \times \left[4\hat{m}_b^2 C_7^{abs}\text{Im}C_7^0 + sC_9^{abs}\text{Im}C_9^0 \right] - 6\sqrt{r}s \left(1 + 2\frac{m_l^2}{q^2} \right) C_9^{abs}\text{Im}C_9^0 \Big\}, \end{aligned} \quad (45)$$

$$\Delta_{FB} = \frac{3}{2R_{\Lambda_b}(s)} \sqrt{1 - \frac{4m_l^2}{q^2}} \phi(s) \text{Im}C_{10} \left[s \left(1 - 2\sqrt{r}\rho - (1-r)\rho^2 \right) C_9^{abs} + 2\hat{m}_b \left(1 - s\rho^2 \right) C_7^{abs} \right], \quad (46)$$

$$\Delta_{PL} = -\frac{1}{R_{\Lambda_b}(s)} \sqrt{1 - \frac{4m_l^2}{q^2}} \text{Im}C_{10} \left\{ C_9^{abs} \left[6\sqrt{r}\rho s (1-r+s) + (1+2s\rho^2)(1-r)^2 + s(1-s\rho^2)(1+r) - s^2(2+s\rho^2) \right] + 6\hat{m}_b C_7^{abs} \left[(1-r-s)(1+s\rho^2) + 4\sqrt{r}\rho s \right] \right\}, \quad (47)$$

$$\Delta_{PT} = \frac{3\pi\hat{m}_l}{4R_{\Lambda_b}(s)} \sqrt{1 - \frac{4m_l^2}{q^2}} \sqrt{s\phi(s)} \left\{ \left(\text{Im}C_9^0 C_{10}^* + C_9^{abs} \text{Im}C_{10} \right) \times \left(1 - 2\sqrt{r}\rho - (s+2t-2r)\rho^2 \right) + \frac{2\hat{m}_b}{s} \left(1 - s\rho^2 \right) \left(\text{Im}C_7^0 C_{10}^* + C_7^{abs} \text{Im}C_{10} \right) \right\}. \quad (48)$$

As seen from the above equations, Δ_Γ is related to $\text{Im}C_7$ and $\text{Im}C_9$, while Δ_{FB} , Δ_{PL} and Δ_{PT} depend on $\text{Im}C_{10}$. Moreover, for small values of $C_9^{abs} \text{Im}C_{10}$ and $C_7^{abs} \text{Im}C_{10}$, Δ_{PT} would still be sizable because $\text{Im}C_9^0 C_{10}^*$ or $\text{Im}C_7^0 C_{10}^*$ would be large.

5 Numerical analysis

In our numerical calculations, the Wilson coefficients are evaluated at the scale $\mu \simeq m_b$ and the other parameters are listed in Table 1 of Ref. [5]. From Eq. (13), we know that the main effects to the deviation of the HQET are from δ . By using a proper nonzero value of δ , we will see later that the deviations of the decay branching ratios of $\Lambda_b \rightarrow \Lambda l^+ l^-$ are only a few percent. Since there is no complete calculation for the form factors of the $\Lambda_b \rightarrow \Lambda$ transition in the literature, we use the form factors derived from QCD sum rule under the assumption of the HQET, given by

$$F_i(q^2) = \frac{F_i(0)}{1 + aq^2 + bq^4}, \quad (49)$$

with the parameters shown in Table 1 of Ref. [6]. In order to illustrate the contributions of new physics, we adopt the results of the generic supersymmetric extension of the SM [22] in which

$$\begin{aligned} C_7^{USY} &= -1.75 (\delta_{23}^u)_{LL} - 0.25 (\delta_{23}^u)_{LR} - 10.3 (\delta_{23}^d)_{LR}, \\ C_9^{USY} &= 0.82 (\delta_{23}^u)_{LR}, \\ C_{10}^{USY} &= -9.37 (\delta_{23}^u)_{LR} + 1.4 (\delta_{23}^u)_{LR} (\delta_{33}^u)_{RL} + 2.7 (\delta_{23}^u)_{LL}, \end{aligned} \quad (50)$$

and take the following values instead of scanning the whole allowed parameter space:

$$\begin{aligned} (\delta_{23}^u)_{LL} &\sim 0.1, \\ (\delta_{33}^u)_{RL} &\sim 0.65, \\ (\delta_{23}^d)_{LR} &\sim 3 \times 10^{-2} e^{i\frac{2\pi}{5}}, \\ (\delta_{23}^u)_{LR} &\sim -0.8 e^{i\frac{\pi}{4}}, \end{aligned} \quad (51)$$

where $(\delta_{ij}^q)_{AB}$ ($i, j = 1, 2, 3$ and $A, B = L, R$) denote the parameters in the mass insertion method, which describe the effects of the flavour violation. The set of the parameters in Eq. (51) satisfies with the constraint from $B \rightarrow X_s \gamma$ on $C_7 = C_7^{SM} + C_7^{SUSY}$ [22]. Hence, the numerical values of the SUSY contributions to the relevant Wilson coefficients are as follows:

$$\begin{aligned} \text{Re } C_7^{SUSY} &\simeq 0.06, & \text{Im } C_7^{SUSY} &\simeq -0.29, \\ \text{Re } C_9^{SUSY} &\simeq -0.46, & \text{Im } C_9^{SUSY} &\simeq -0.46, \\ \text{Re } C_{10}^{SUSY} &\simeq 4.78, & \text{Im } C_{10}^{SUSY} &\simeq 4.50. \end{aligned} \quad (52)$$

We note that the contributions of the minimal supersymmetric standard model (MSSM) to $b \rightarrow sl^+l^-$ can be found in Refs. [23] and [24].

To show the typical values of various asymmetries, we define the integrated quantities as

$$\bar{Q} = \int_{s_{\min}}^{s_{\max}} Q(s) ds \quad (53)$$

where Q denote the physical observables with $s_{\min} = 4\hat{m}_l$ and $s_{\max} = (1 - \sqrt{r})^2$.

5.1 Decay rates and invariant mass distributions

We now discuss the influences of δ , ρ , and ω on the branching ratios (BRs) of $\Lambda_b \rightarrow \Lambda l^+ l^-$ decays in detail. The effects of k_j for compensating the assumption of the FA and VMD have been analyzed in [5]. In Table 2, we show the BRs by choosing different sets of parameters. Our results are given as follows:

Table 2: BRs (in the unit of 10^{-6}) for various parameters with $\omega = 0$ and neglecting LD effects.

Parameter	$\Lambda_b \rightarrow \Lambda e^+ e^-$	$\Lambda_b \rightarrow \Lambda \mu^+ \mu^-$	$\Lambda_b \rightarrow \Lambda \tau^+ \tau^-$
HQET	2.23	2.08	1.79×10^{-1}
$\delta = 0.05$	2.36	2.21	1.86×10^{-1}
$\delta = -0.05$	2.09	1.96	1.71×10^{-1}
$\rho = 0, \delta = 0$	2.52	2.38	2.66×10^{-1}
$C_7 = 0, \delta = 0$	2.36	2.34	2.23×10^{-1}
$C_7 = -C_7^{SM}, \delta = 0$	3.34	3.19	2.76×10^{-1}

1. By taking $|\delta| = 0.05$ which means 10% away from that in to the HQET, we clearly see that the deviations of the BRs are only 4 – 6%. It is a good approximation to neglect the explicit δ term in Eqs. (17), (22) and (45). Hence, $\bar{f} = (f_1 + g_1)/2$, which also owns the δ effect, is the main nonperturbative part.
2. If $\rho = 0$, the effects are about 10% for e and μ modes but 48% for τ one.
3. If one neglects the contribution from C_7 , the influences on $B(\Lambda_b \rightarrow \Lambda l^+ l^-)$ for e , μ and τ modes are about 5%, 12% and 24%, respectively. However, taking the magnitude of C_7 is the same as the SM but with an opposite sign, the deviations are all over 50%.

The contributions of the parameter ω to the BRs of $\Lambda_b \rightarrow \Lambda l^+ l^-$ are listed in Table 3 and the invariant mass distributions are shown in Figure 1. From Table 3, it is clear that the nonfactorizable effects are small on the BRs. However, those directly related to ω effects such as P_T and CP asymmetries will have large influences.

Table 3: BRs (in the unit of 10^{-6}) without LD effects with different values of ω

Mode	$\Lambda_b \rightarrow \Lambda e^+ e^-$	$\Lambda_b \rightarrow \Lambda \mu^+ \mu^-$	$\Lambda_b \rightarrow \Lambda \tau^+ \tau^-$
$\omega = 0.15$	2.24	2.12	1.89×10^{-1}
$\omega = 0.$	2.23	2.08	1.79×10^{-1}
$\omega = -0.15$	2.25	2.06	1.71×10^{-1}

As for the new physics contributions, using the values of the SUSY model in Eq. (52), we show the results in Table 4. Although the deviations of the BRs to the SM are not significant, they have a large effect on the lepton and CP asymmetries which will be shown next.

Table 4: BRs (in unit of 10^{-6}) in the generic SUSY model.

Model	$\Lambda_b \rightarrow \Lambda e^+ e^-$	$\Lambda_b \rightarrow \Lambda \mu^+ \mu^-$	$\Lambda_b \rightarrow \Lambda \tau^+ \tau^-$
SUSY	2.47	2.24	1.79×10^{-1}

5.2 Forward-backward and lepton polarization asymmetries

From Eq. (14) and $R = F_2/F_1 \simeq -0.25$ in the HQET, we have that $\rho \simeq -0.26$. We note that ρ is defined by the ratio of the form factors and it is expected to be insensitive to the QCD models. With $s_{\max} \simeq 0.64$, we obtain $s_{\max} \rho^2 \simeq 0.04$, $(1-r)\rho^2 \simeq 0.06$ and $2\sqrt{r}\rho \simeq 0.2$. Using these values, one can simplify Eqs. (30) and (31) to

$$R_{FB}(s) \simeq s \left(1 - 2\sqrt{r}\rho \right) \text{Re}C_9^{\text{eff}} C_{10}^* + 2\hat{m}_b \text{Re}C_7^{\text{eff}} C_{10}^* \quad (54)$$

and

$$\text{Re}C_9^{\text{eff}} C_{10}^* \simeq -\frac{2\hat{m}_b}{s_0(1-2\sqrt{r}\rho)} \text{Re}C_7^{\text{eff}} C_{10}^*, \quad (55)$$

respectively. It is easy to see that s_0 is only sensitive to the Wilson coefficients. The result is similar to the case in $B \rightarrow K^* l^+ l^-$ decay [3, 19] where the approximation of the large energy effective theory (LEET) [25] is used. As for the lepton polarization asymmetries, with the same approximation, Eqs. (37) and (38) can also be reduced to

$$R_L \simeq -\text{Re}C_9^{\text{eff}} C_{10}^* \left[1 + s - 2s^2 + 6\sqrt{r}\rho s(1+s) \right] - 6\hat{m}_b \text{Re}C_7^{\text{eff}} C_{10}^* \left[1 - s + 4\sqrt{r}\rho s \right], \quad (56)$$

$$R_T \simeq \left(1 - 2\sqrt{r}\rho \right) \text{Im}C_9^{\text{eff}} C_{10}^* + \frac{2\hat{m}_b}{s} \text{Im}C_7^{\text{eff}} C_{10}^*, \quad (57)$$

respectively. Hence, the lepton asymmetries are all more sensitive to the Wilson coefficients than the nonperturbative QCD effects.

It is worth to mention that the effects of ω , introduced for the LD contributions to $b \rightarrow s\gamma$ and absorbed to C_7^{eff} , will change $\text{Re}C_7^{eff}$ in the SM such that s_0 is also shifted. Therefore, in terms of s_0 , we can also theoretically determine ω by comparing the result with that of $\omega = 0$. Another interesting quantity is T-odd observable of P_T which is proportional to $C_{10}\text{Im}C_7^{eff}$ in the SM. Due to the enhancement of C_{10} , a nonzero value of ω will modify P_T enormously. As for the other asymmetries, the effects are insignificant. The estimations of integrated lepton asymmetries with different values of ω in the SM are displayed in Table 5 and the corresponding distributions are shown in Figures 2 – 4.

Table 5: Integrated lepton asymmetries in the SM without LD effects.

Parameter	Mode	$10^2 \bar{A}_{FB}$	$10^2 \bar{P}_L$	$10^2 \bar{P}_T$
$\omega = 0.15$	$\Lambda_b \rightarrow \Lambda \mu^+ \mu^-$	-14.37	59.50	0.11
	$\Lambda_b \rightarrow \Lambda \tau^+ \tau^-$	-3.98	10.70	0.53
$\omega = 0.$	$\Lambda_b \rightarrow \Lambda \mu^+ \mu^-$	-13.38	58.30	0.07
	$\Lambda_b \rightarrow \Lambda \tau^+ \tau^-$	-3.99	10.84	0.39
$\omega = -0.15$	$\Lambda_b \rightarrow \Lambda \mu^+ \mu^-$	-12.24	56.70	0.04
	$\Lambda_b \rightarrow \Lambda \tau^+ \tau^-$	-4.00	10.94	0.23

To illustrate the new physics effects, the integrated lepton asymmetries in the generic SUSY model with ω are listed in Table 6 and their distributions as a function of q^2/M_{Λ_b} are shown in Figures 5 – 7. From the figures, we see that SUSY effects make the shapes of lepton asymmetries quite differ from that in the SM. We summary the results as follows:

1. Since the SD contributions to $C_9 C_{10}^*$ and $\text{Re}C_7 C_{10}^*$ are -1.40 and -1.35, respectively, which violate the condition in Eq. (55), the vanishing point is removed.
2. Due to the factor of \hat{m}_b/s , from Figure 7, we see that $\text{Im}C_7^{eff} C_{10}^*$ has a large effect on P_T in the small s region.
3. In the SUSY model, P_T could reach 1% and 10% for the light lepton and τ modes, which are only 0.2% and 3% at most in the SM, respectively.

Table 6: Integrated lepton asymmetries in the generic SUSY model with $\omega = 0$.

Mode	$10^2 \bar{A}_{FB}$	$10^2 \bar{P}_L$	$10^2 \bar{P}_T$
$\Lambda_b \rightarrow \Lambda \mu^+ \mu^-$	-10.53	24.46	-0.57
$\Lambda_b \rightarrow \Lambda \tau^+ \tau^-$	-1.84	4.40	-2.51

Table 7: CP asymmetries in the generic SUSY model for different values of ω

Parameter	Mode	$10^2 \bar{\Delta}_\Gamma$	$10^2 \bar{\Delta}_{FBA}$	$10^2 \bar{\Delta}_{P_L}$	$10^2 \bar{\Delta}_{P_T}$
$\omega = 0.15$	$\Lambda_b \rightarrow \Lambda \mu^+ \mu^-$	2.05	-2.62	6.48	-0.53
	$\Lambda_b \rightarrow \Lambda \tau^+ \tau^-$	1.83	-0.79	1.94	-2.01
$\omega = 0.$	$\Lambda_b \rightarrow \Lambda \mu^+ \mu^-$	1.59	-1.89	5.00	-0.47
	$\Lambda_b \rightarrow \Lambda \tau^+ \tau^-$	1.38	-0.59	1.53	-2.21
$\omega = -0.15$	$\Lambda_b \rightarrow \Lambda \mu^+ \mu^-$	1.05	-1.06	3.34	-0.40
	$\Lambda_b \rightarrow \Lambda \tau^+ \tau^-$	0.89	-0.37	1.06	-2.41

5.3 CP asymmetries

In the SM, for $b \rightarrow sl^+l^-$, the relevant CKM matrix element is $V_{tb}V_{ts}^*$ which is real under the Wolfenstein's parametrization. Nonzero CPAs will indicate clearly the existence of new physics. We remark that the CPAs can be in fact induced by the complex CKM matrix element $V_{ub}V_{us}^*$ which is also the source of the direct CPA in $B \rightarrow X_s \gamma$ in the SM. However, we expect that such effects to the CPAs in $b \rightarrow sl^+l^-$ are smaller than that in $B \rightarrow X_s \gamma$ where the CPA is less than 1%. The main reason for the smallness is because of the presences of C_9 and C_{10} contributions to the rates of $b \rightarrow sl^+l^-$, which are absent in $B \rightarrow X_s \gamma$.

With the values in Eq. (52), the averaged CPAs in the generic SUSY model for $\Lambda_b \rightarrow \Lambda l^+l^-$ are listed in Table 7 and their distributions as a function of $s = q^2/M_{\Lambda_b}^2$ are shown in Figures 8 – 11. The results are given as follows:

1. From Eqs. (45) and (48), we see that the terms corresponding to $C_7^{abs} Im C_7^0$ and $Im C_7^0 C_{10}^* + C_7^{abs} Im C_{10}$ are associated with a factor of \hat{m}_b/s . If sizable imaginary parts exist, in the small s region the distributions will be significant. Due to this reason, in Figure 8a one finds that $\Delta_\Gamma(s)$ for $\Lambda_b \rightarrow \Lambda l^+l^-$ ($l = e, \mu$) increase as s decreases. On the other hand, if the term with \hat{m}_b/s in Eq. (48) is dropped, the distributions of $\Delta_{P_T}(s)$ for e and μ modes do not contain zero value. We note that with the values in Eq. (52), the main effect on $\Delta_\Gamma(s)$ in the small s region is from $C'_{b \rightarrow s \gamma}$.
2. $\Delta_{P_L}(s)$ for all lepton channels and $\Delta_{P_T}(s)$ for the τ one could be over 10%, while the remaining CP asymmetries are at the level of a few percent. We remark that if we can scan all the allowed SUSY parameters, the asymmetries except $\Delta_{P_T}(s)$ for lighter lepton modes would reach up 10%.
3. It is known that $\Delta_{P_T}(s)$ is a T-odd observable and the other CPAs belong to the direct CP violation which needs absorptive parts in the processes. This is the reason why the distributions of $\Delta_\Gamma(s)$, $\Delta_{FBA}(s)$ and $\Delta_{P_L}(s)$ around the RE region have the similar shapes but are different from that of $\Delta_{P_T}(s)$. Moreover, all the direct CPAs are sensitive to ω unlike the cases of the CP conserving lepton asymmetries discussed in Sec. 5.2.

6 Conclusions

We have given a systematic study on the rare baryonic decays of $\Lambda_b \rightarrow \Lambda l^+ l^-$ ($l = e, \mu, \tau$). For the $\Lambda_b \rightarrow \Lambda$ transition, we have related all the form factors with F_1 and F_2 , and we have found that $\delta = 0$ and $\rho \simeq R \equiv F_2/F_1$, in the limit of the HQET. Inspired by the HQEF, we have presented the differential decay rates and the di-lepton forward-backward, lepton polarization and four possible CP violating asymmetries in terms of the parameters \bar{f} , δ and ρ . We have shown that the non-factorizable effects for the BRs and CP-even lepton asymmetries are small but large for P_T and the direct CPAs. We have also demonstrated that most of the observables such as \mathcal{A}_{FB} , $P_{L,T}$ and Δ_α ($\alpha = \Gamma, FB, P_L$ and P_T), are insensitive to the non-perturbative QCD effects. We have illustrated our results in the specific CP violating SUSY model. We have found that all the direct CP violating asymmetries are in the level of 1 – 10%. To measure these asymmetries at the $n\sigma$ level, for example, in the tau mode, at least $0.5n^2 \times (10^9 - 10^{10})$ Λ_b decays are required. It could be done in the second generation B-physics experiments, such as LHCb, ATLAS, and CMS at the LHC, and BTeV at the Tevatron, which produce $\sim 10^{12} b\bar{b}$ pairs per year [26] Finally we remark that measuring these CPAs at a level of 10^{-2} is a clear indication of new CP violation mechanism beyond the SM.

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Appendix

A. Form factors and decay amplitudes

For the exclusive decays involving $\Lambda_b(p_{\Lambda_b}) \rightarrow \Lambda(p_\Lambda)$, the transition form factors can be parametrized generally as follows:

$$\langle \Lambda | \bar{s} \gamma_\mu b | \Lambda_b \rangle = f_1 \bar{u}_\Lambda \gamma_\mu u_{\Lambda_b} + f_2 \bar{u}_\Lambda i \sigma_{\mu\nu} q^\nu u_{\Lambda_b} + f_3 q_\mu \bar{u}_\Lambda u_{\Lambda_b}, \quad (58)$$

$$\langle \Lambda | \bar{s} \gamma_\mu \gamma_5 b | \Lambda_b \rangle = g_1 \bar{u}_\Lambda \gamma_\mu \gamma_5 u_{\Lambda_b} + g_2 \bar{u}_\Lambda i \sigma_{\mu\nu} q^\nu \gamma_5 u_{\Lambda_b} + g_3 q_\mu \bar{u}_\Lambda \gamma_5 u_{\Lambda_b}, \quad (59)$$

$$\begin{aligned} \langle \Lambda | \bar{s} i \sigma_{\mu\nu} b | \Lambda_b \rangle &= f_T \bar{u}_\Lambda i \sigma_{\mu\nu} u_{\Lambda_b} + f_T^V \bar{u}_\Lambda (\gamma_\mu q_\nu - \gamma_\nu q_\mu) u_{\Lambda_b} \\ &\quad + f_T^S (P_\mu q_\nu - P_\nu q_\mu) \bar{u}_\Lambda u_{\Lambda_b}, \end{aligned} \quad (60)$$

$$\begin{aligned} \langle \Lambda | \bar{s} i \sigma_{\mu\nu} \gamma_5 b | \Lambda_b \rangle &= g_T \bar{u}_\Lambda i \sigma_{\mu\nu} \gamma_5 u_{\Lambda_b} + g_T^V \bar{u}_\Lambda (\gamma_\mu q_\nu - \gamma_\nu q_\mu) \gamma_5 u_{\Lambda_b} \\ &\quad + g_T^S (P_\mu q_\nu - P_\nu q_\mu) \bar{u}_\Lambda \gamma_5 u_{\Lambda_b}, \end{aligned} \quad (61)$$

where $P = p_{\Lambda_b} + p_\Lambda$, $q = p_{\Lambda_b} - p_\Lambda$ and form factors, $\{f_i\}$ and $\{g_i\}$, are all functions of q^2 . Using the equations of the motion, we have

$$(M_\Lambda + M_{\Lambda_b}) \bar{u}_\Lambda \gamma_\mu u_{\Lambda_b} = (p_{\Lambda_b} + p_\Lambda)_\mu \bar{u}_\Lambda u_{\Lambda_b} + i \bar{u}_\Lambda \sigma_{\mu\nu} q^\nu u_{\Lambda_b}, \quad (62)$$

$$(M_\Lambda - M_{\Lambda_b}) \bar{u}_\Lambda \gamma_\mu \gamma_5 u_{\Lambda_b} = (p_{\Lambda_b} + p_\Lambda)_\mu \bar{u}_\Lambda \gamma_5 u_{\Lambda_b} + i \bar{u}_\Lambda \sigma_{\mu\nu} q^\nu \gamma_5 u_{\Lambda_b}. \quad (63)$$

The form factors for dipole operators are derived as

$$\langle \Lambda | \bar{s} i \sigma_{\mu\nu} q^\nu b | \Lambda_b \rangle = f_1^T \bar{u}_\Lambda \gamma_\mu u_{\Lambda_b} + f_2^T \bar{u}_\Lambda i \sigma_{\mu\nu} q^\nu u_{\Lambda_b} + f_3^T q_\mu \bar{u}_\Lambda u_{\Lambda_b}, \quad (64)$$

$$\langle \Lambda | \bar{s} i \sigma_{\mu\nu} q^\nu \gamma_5 b | \Lambda_b \rangle = g_1^T \bar{u}_\Lambda \gamma_\mu \gamma_5 u_{\Lambda_b} + g_2^T \bar{u}_\Lambda i \sigma_{\mu\nu} q^\nu \gamma_5 u_{\Lambda_b} + g_3^T q_\mu \bar{u}_\Lambda \gamma_5 u_{\Lambda_b}. \quad (65)$$

with

$$\begin{aligned} f_2^T &= f_T - f_T^S q^2, \\ f_1^T &= [f_T^V + f_T^S (M_\Lambda + M_{\Lambda_b})] q^2, \\ f_1^T &= -\frac{q^2}{(M_{\Lambda_b} - M_\Lambda)} f_3^T, \\ g_2^T &= g_T - g_T^S q^2, \\ g_1^T &= [g_T^V + g_T^S (M_\Lambda - M_{\Lambda_b})] q^2, \\ g_1^T &= \frac{q^2}{(M_{\Lambda_b} + M_\Lambda)} g_3^T. \end{aligned} \quad (66)$$

We now give the most general formulas by including the right-handed coupling in the effective Hamiltonian with a complete set of form factors. The free quark decay amplitudes for $b \rightarrow sl^+ l^-$ are given by

$$\begin{aligned} \mathcal{H}(b \rightarrow sl^+ l^-) &= \frac{G_F \alpha_{em}}{\sqrt{2} \pi} V_{tb} V_{ts}^* \left[\bar{s} \gamma^\mu (C_9^L P_L + C_9^R P_R) b \bar{l} \gamma_\mu l \right. \\ &\quad + \bar{s} \gamma^\mu (C_{10}^L P_L + C_{10}^R P_R) b \bar{l} \gamma_\mu \gamma_5 l \\ &\quad \left. - \frac{2m_b}{q^2} \bar{s} i \sigma_{\mu\nu} q^\nu (C_7^L P_R + C_7^R P_L) b \bar{l} \gamma_\mu l \right] \end{aligned} \quad (67)$$

where C_i^L and C_i^R ($i = 7, 9, 10$) denote the effective Wilson coefficients of left- and right-handed couplings, respectively. With the most general form factors in Eqs. (58), (59), (64) and (65), and the effective Hamiltonian in Eq. (67), the transition matrix elements for the decays of $\Lambda_b \rightarrow \Lambda l^+ l^-$ are expressed as

$$\begin{aligned} \mathcal{M}(\Lambda_b \rightarrow \Lambda l^+ l^-) = & \frac{G_F \alpha_{em}}{\sqrt{2}} V_{tb} V_{ts}^* \left\{ \left[\bar{\Lambda} \gamma_\mu (A_1 P_R + B_1 P_L) \Lambda_b \right. \right. \\ & + \bar{\Lambda} i \sigma_{\mu\nu} q^\nu (A_2 P_R + B_2 P_L) \Lambda_b \left. \right] \bar{l} \gamma_\mu l \\ & + \left[\bar{\Lambda} \gamma_\mu (D_1 P_R + E_1 P_L) \Lambda_b + \bar{\Lambda} i \sigma_{\mu\nu} q^\nu (D_2 P_R + E_2 P_L) \Lambda_b \right. \\ & \left. \left. + q_\mu \bar{\Lambda} (D_3 P_R + E_3 P_L) \Lambda_b \right] \bar{l} \gamma_\mu \gamma_5 l \right\} \end{aligned} \quad (68)$$

where

$$\begin{aligned} A_i &= C_9^R \frac{f_i + g_i}{2} - \frac{2m_b}{q^2} C_7^R \frac{f_i^T - g_i^T}{2} + C_9^L \frac{f_i - g_i}{2} - \frac{2m_b}{q^2} C_7^L \frac{f_i^T + g_i^T}{2}, \\ B_i &= C_9^L \frac{f_i + g_i}{2} - \frac{2m_b}{q^2} C_7^L \frac{f_i^T - g_i^T}{2} + C_9^R \frac{f_i - g_i}{2} - \frac{2m_b}{q^2} C_7^R \frac{f_i^T + g_i^T}{2}, \\ D_i &= C_{10}^R \frac{f_i + g_i}{2} + C_{10}^L \frac{f_i - g_i}{2}, \\ E_i &= C_{10}^L \frac{f_i + g_i}{2} + C_{10}^R \frac{f_i - g_i}{2}. \end{aligned} \quad (69)$$

B. Differential decay rates

Using the transition matrix elements in Eq. (68), the double differential decay rates can be derived as

$$\frac{d\Gamma}{ds d\hat{z}} = \frac{G_F^2 \alpha_{em}^2 \lambda_i^2}{768 \pi^5} M_{\Lambda_b}^5 \sqrt{\phi(s)} \sqrt{1 - \frac{4m_l^2}{q^2}} \bar{f}^2 R_{\Lambda_b}(s, \hat{z}), \quad (70)$$

where

$$R_{\Lambda_b}(s, \hat{z}) = I_0(s, \hat{z}) + \hat{z} I_1(s, \hat{z}) + \hat{z}^2 I_2(s, \hat{z}) \quad (71)$$

with

$$\begin{aligned} I_0(s, \hat{z}) = & -6\sqrt{r}\hat{s} \left[\left(1 + 2\frac{m_l^2}{q^2} \right) \text{Re} A_1 B_1^* + \left(1 - 6\frac{m_l^2}{q^2} \right) \text{Re} D_1 E_1^* \right] \\ & + \frac{3}{4} \left((1-r)^2 - s^2 \right) \left(|A_1|^2 + |B_1|^2 + |D_1|^2 + |E_1|^2 \right) \\ & + 6\hat{m}_l^2 t \left(|A_1|^2 + |B_1|^2 - |D_1|^2 - |E_1|^2 \right) \\ & + 12\hat{m}_l^2 M_{\Lambda_b} \sqrt{r} (1-t) (\text{Re} D_1 D_3'^* + \text{Re} E_1 E_3'^*) \\ & + 12\hat{m}_l^2 M_{\Lambda_b} (t-r) (\text{Re} D_1 E_3'^* + \text{Re} D_3 E_1'^*) \end{aligned}$$

$$\begin{aligned}
& +6M_{\Lambda_b}\sqrt{r}s(1-t)\left[\left(1+2\frac{m_l^2}{q^2}\right)(\text{Re}A_1A_2^*+\text{Re}B_1B_2^*)\right. \\
& \left.+\left(1-2\frac{m_l^2}{q^2}\right)(\text{Re}D_1D_2^*+\text{Re}E_1E_2^*)\right] \\
& -6M_{\Lambda_b}s(t-r)\left[\left(1+2\frac{m_l^2}{q^2}\right)(\text{Re}A_1B_2^*+\text{Re}A_2B_1^*)\right. \\
& \left.+\left(1-6\frac{m_l^2}{q^2}\right)(\text{Re}D_1E_2^*+\text{Re}D_2E_1^*)\right] \\
& -6M_{\Lambda_b}^2\sqrt{r}s^2\left(1+2\frac{m_l^2}{q^2}\right)\text{Re}A_2B_2^*-6M_{\Lambda_b}^2\sqrt{r}s^2\left(1-6\frac{m_l^2}{q^2}\right)\text{Re}D_2E_2^* \\
& -6M_{\Lambda_b}^2\left[s(1-t)(t-r)-\frac{1}{8}\left((1-r)^2-s^2\right)\right]\left(|A_2|^2+|B_2|^2+|D_2|^2+|E_2|^2\right) \\
& -6M_{\Lambda_b}^2\hat{m}_l^2(2r-(1+r)t)\left(|A_2|^2+|B_2|^2-|D_2|^2-|E_2|^2\right) \\
& +12\hat{m}_l^2M_{\Lambda_b}^2st(\text{Re}D_2D_3'^*+\text{Re}E_2E_3'^*)+12\hat{m}_l^2M_{\Lambda_b}^2\sqrt{r}s(\text{Re}D_2E_3'^*+\text{Re}D_3'E_2) , \\
I_1(s,\hat{z}) = & 3s\phi(s)\left\{-\left(\text{Re}A_1D_1^*-\text{Re}B_1E_1^*\right)+M_{\Lambda_b}\left[\sqrt{r}(\text{Re}A_1D_2^*-\text{Re}B_1E_2^*)\right.\right. \\
& \left.+\left(\text{Re}A_1E_2^*-\text{Re}B_1D_2^*\right)+\sqrt{r}(\text{Re}A_2D_1^*-\text{Re}B_2E_1^*)-\left(\text{Re}A_2E_1^*-\text{Re}B_2D_1^*\right)\right] \\
& \left.+M_{\Lambda_b}(1-r)(\text{Re}A_2D_2^*-\text{Re}B_2E_2^*)\right\} , \\
I_2(s,\hat{z}) = & -\frac{3}{4}\phi(s)\left(1-4\frac{m_l^2}{q^2}\right)\left(|A_1|^2+|B_1|^2+|D_1|^2+|E_1|^2\right) \\
& +\frac{3}{4}M_{\Lambda_b}^2\phi(s)\left(1-4\frac{m_l^2}{q^2}\right)\left(|A_2|^2+|B_2|^2+|D_2|^2+|E_2|^2\right) \tag{72}
\end{aligned}$$

where $D_3' = D_3 - D_2$ and $E_3' = E_3 - E_2$, and $\hat{z} = \hat{p}_B \cdot \hat{p}_{l^+}$ denotes the angle between the momentum of Λ_b and that of l^+ in the di-lepton invariant mass frame.

C. Forward-backward and lepton asymmetries

From Eq. (68), the functions of R_{Λ_b} and R_{FB} in the differential and normalized FBAs for $\Lambda_b \rightarrow \Lambda l^+ l^-$ in Eqs. (28) and (29) are given by

$$R_{\Lambda_b}(s) = \frac{1}{2} \int_{-1}^1 d\hat{z} R_{\Lambda_b}(s, \hat{z}) \tag{73}$$

and

$$\begin{aligned}
R_{FB}(s) = & -\text{Re}(A_1D_1^* - B_1E_1^*) + M_{\Lambda_b}\left[\sqrt{r}\text{Re}(A_1D_2^* - B_1E_2^*) + \text{Re}(A_1E_2^* - B_1D_2^*)\right. \\
& \left.+\sqrt{r}\text{Re}(A_2D_1^* - B_2E_1^*) - \text{Re}(A_2E_1^* - B_2D_1^*)\right] \\
& +M_{\Lambda_b}^2(1-r)\text{Re}(A_2D_2^* - B_2E_2^*) , \tag{74}
\end{aligned}$$

respectively. We can also define the longitudinal and transverse lepton polarization asymmetries. Explicitly, by the definition of Eq. (34), we get

$$P_L = -\frac{1}{R_{\Lambda_b}(s)}\sqrt{1-\frac{4\hat{m}_l^2}{s}}\left\{\left(s(1+r-s)+(1-r)^2-s^2\right)\text{Re}(A_1D_1^*+B_1E_1^*)\right.$$

$$\begin{aligned}
& +M_{\Lambda_b}^2 s \left(-s(1+r-s) + 2(1-r)^2 - 2s^2 \right) \text{Re} (A_2 D_2^* + B_2 E_2^*) \\
& -6s\sqrt{r} \left[\text{Re} (A_1 E_1^* + B_1 D_1^*) + M_{\Lambda_b}^2 s \text{Re} (A_2 E_2^* + B_2 D_2^*) \right] \\
& +3sM_{\Lambda} (1-r+s) [\text{Re} (A_1 D_2^* + B_1 E_2^*) + \text{Re} (A_2 D_1^* + B_2 E_1^*)] \\
& -3sM_{\Lambda_b} (1-r-s) [\text{Re} (A_1 E_2^* + B_1 D_2^*) - \text{Re} (A_2 E_1^* + B_2 D_1^*)], \tag{75}
\end{aligned}$$

$$\begin{aligned}
P_T = & \frac{3}{4} \pi \hat{m}_l \sqrt{1 - \frac{4m_l^2}{q^2}} \sqrt{s\phi(s)} \frac{1}{R_{\Lambda_b}(s)} \{ -\text{Im} (A_1 D_1^* - B_1 E_1^*) \\
& +M_{\Lambda} [\text{Im} (A_1 D_2^* - B_1 E_2^*) + \text{Im} (A_2 D_1^* - B_2 E_1^*)] \\
& +M_{\Lambda_b} [\text{Im} (A_1 E_2^* - B_1 D_2^*) - \text{Im} (A_2 E_1^* - B_2 D_1^*)] \\
& +M_{\Lambda_b}^2 (1-r) \text{Im} (A_2 D_2^* - B_2 E_2^*) \}. \tag{76}
\end{aligned}$$

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Figure Captions

- Figure 1: BRs as a function of $q^2/M_{\Lambda_b}^2$ for (a) $\Lambda_b \rightarrow \Lambda\mu^+\mu^-$ and (b) $\Lambda_b \rightarrow \Lambda\tau^+\tau^-$. The curves with and without resonant shapes represent including and no LD contributions, respectively. The dashed, solid and dash-dotted curves stand for $\omega = 0.15, 0$, and -0.15 , respectively.
- Figure 2: Same as Figure 1 but for the FBAs.
- Figure 3: Same as Figure 1 but for the longitudinal polarization asymmetries.
- Figure 4: Same as Figure 1 but for the transverse polarization asymmetries.
- Figure 5: FBAs in the generic SUSY model as a function of $q^2/M_{\Lambda_b}^2$ for (a) $\Lambda_b \rightarrow \Lambda\mu^+\mu^-$ and (b) $\Lambda_b \rightarrow \Lambda\tau^+\tau^-$. The solid and dashed curves stand for the SM and SUSY model, respectively.
- Figure 6: Same as Figure 5 but for the longitudinal polarization asymmetries.
- Figure 7: Same as Figure 5 but for the transverse polarization asymmetries.
- Figure 8: Same as Figure 5 but for Δ_Γ .
- Figure 9: Same as Figure 5 but for Δ_{FBA} .
- Figure 10: Same as Figure 5 but for Δ_{P_L} .
- Figure 11: Same as Figure 5 but for Δ_{P_T} .

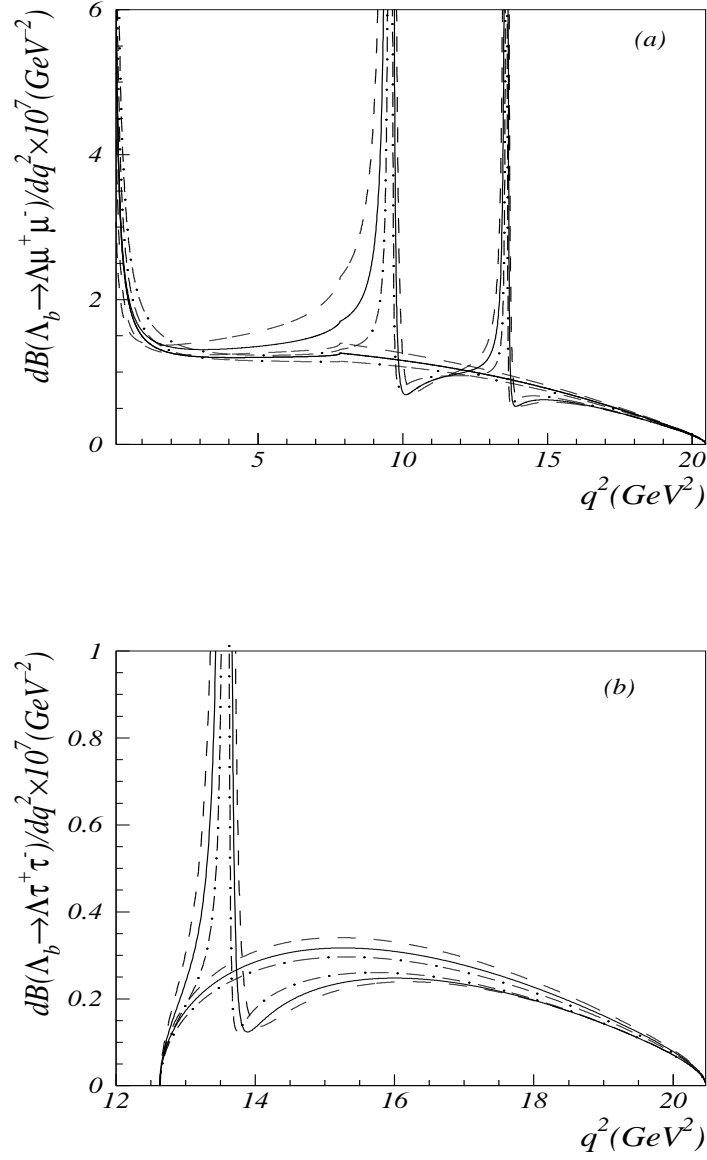


Figure 1: BRs as a function of $q^2/M_{\Lambda_b}^2$ for (a) $\Lambda_b \rightarrow \Lambda \mu^+ \mu^-$ and (b) $\Lambda_b \rightarrow \Lambda \tau^+ \tau^-$. The curves with and without resonant shapes represent including and no LD contributions, respectively. The dashed, solid and dash-dotted curves stand for $\omega = 0.15$, 0, and -0.15 , respectively.

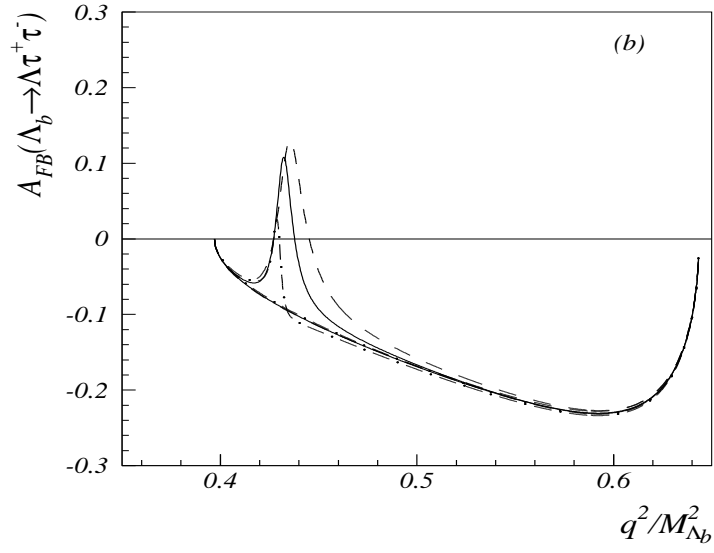
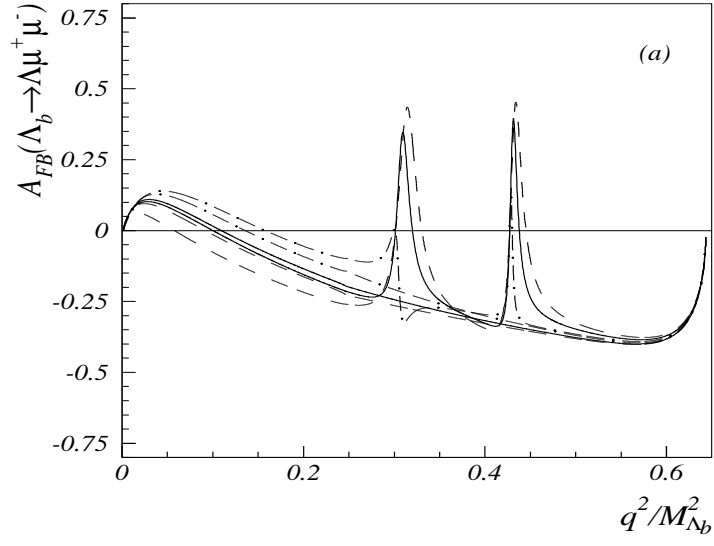


Figure 2: Same as Figure 1 but for the FBAs.

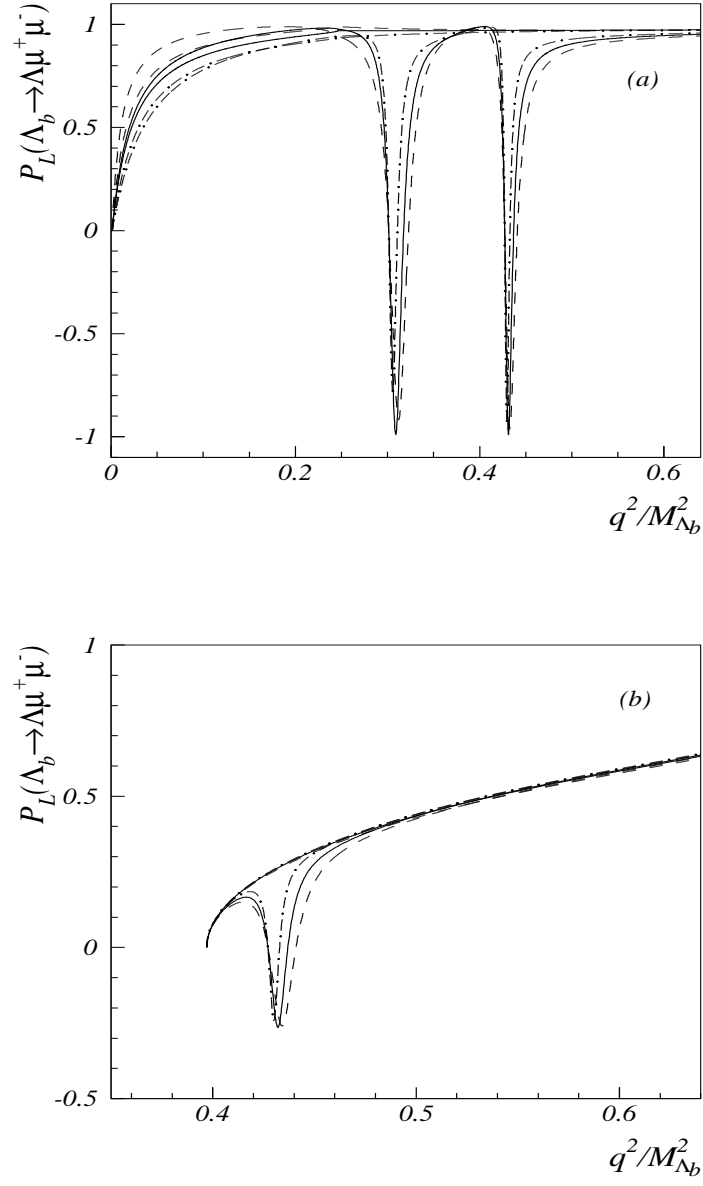


Figure 3: Same as Figure 1 but for the longitudinal polarization asymmetries.

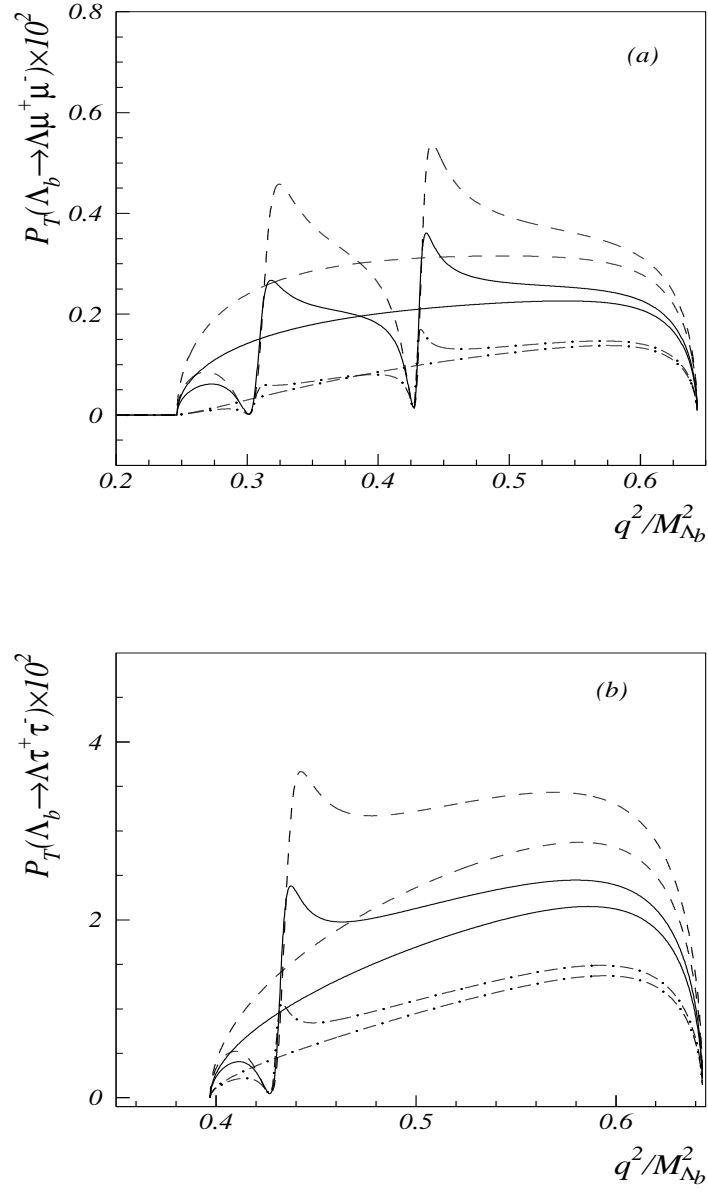


Figure 4: Same as Figure 1 but for the transverse polarization asymmetries.

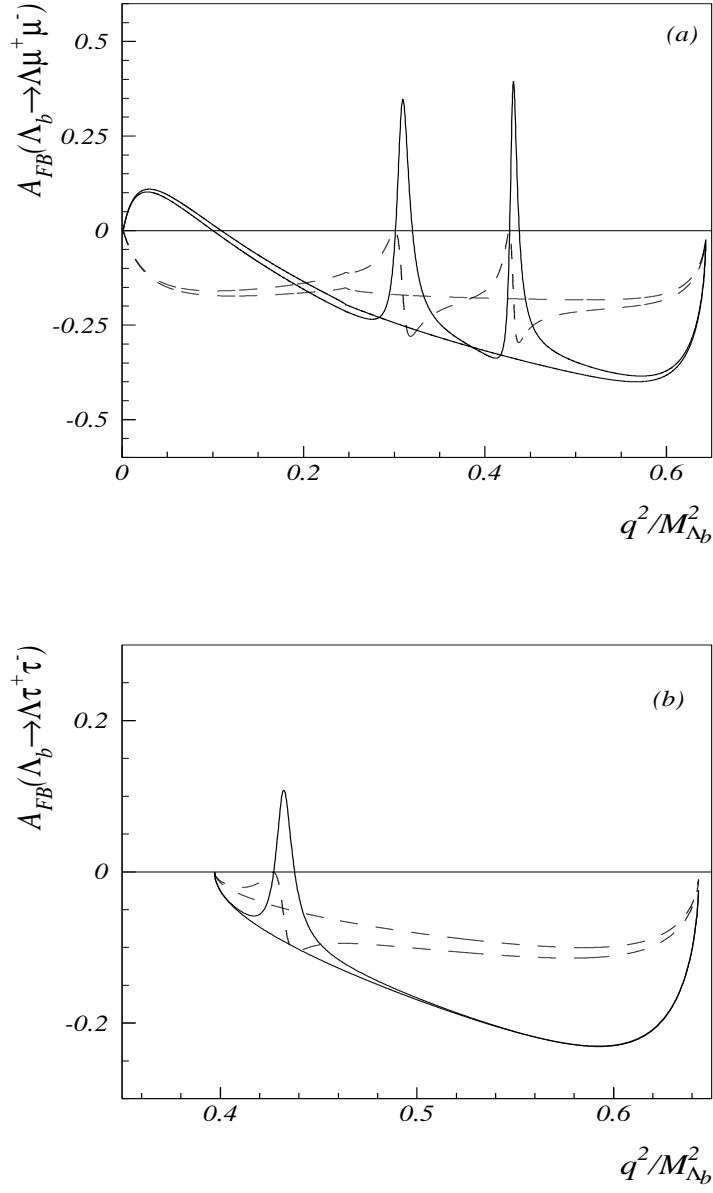


Figure 5: FBAs in the generic SUSY model as a function of $q^2/M_{\Lambda_b}^2$ for (a) $\Lambda_b \rightarrow \Lambda \mu^+ \mu^-$ and (b) $\Lambda_b \rightarrow \Lambda \tau^+ \tau^-$. The solid and dashed curves stand for the SM and SUSY model, respectively.

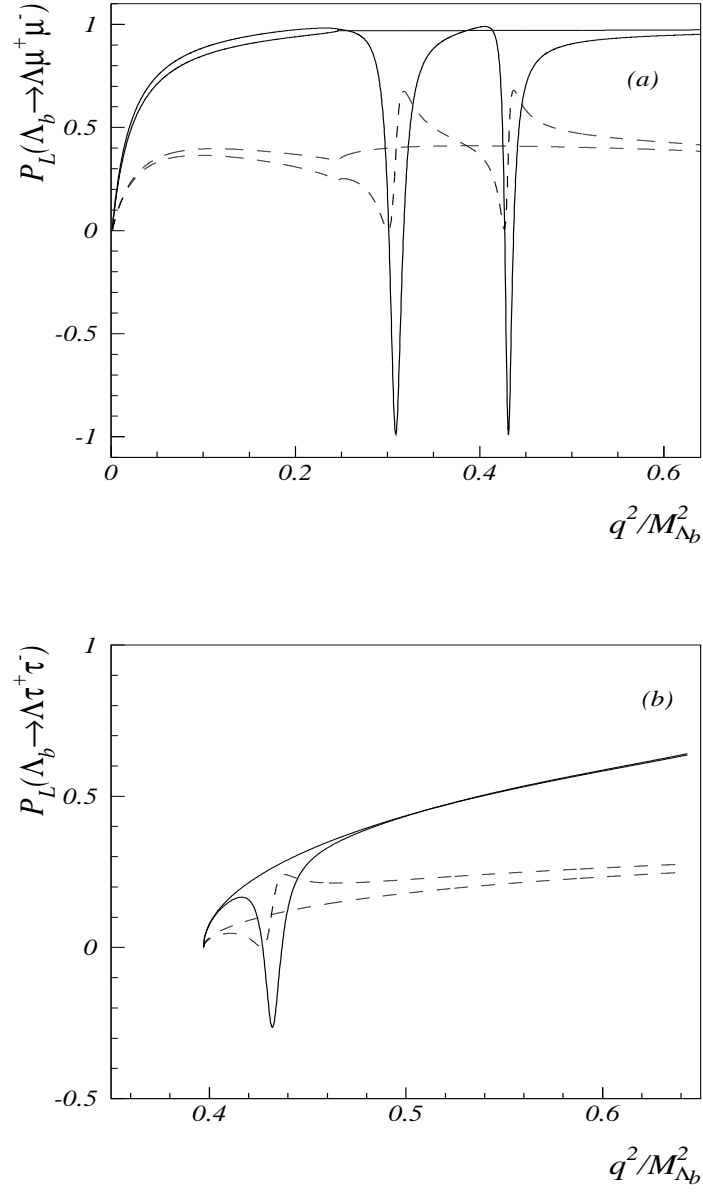


Figure 6: Same as Figure 5 but for the longitudinal polarization asymmetries.

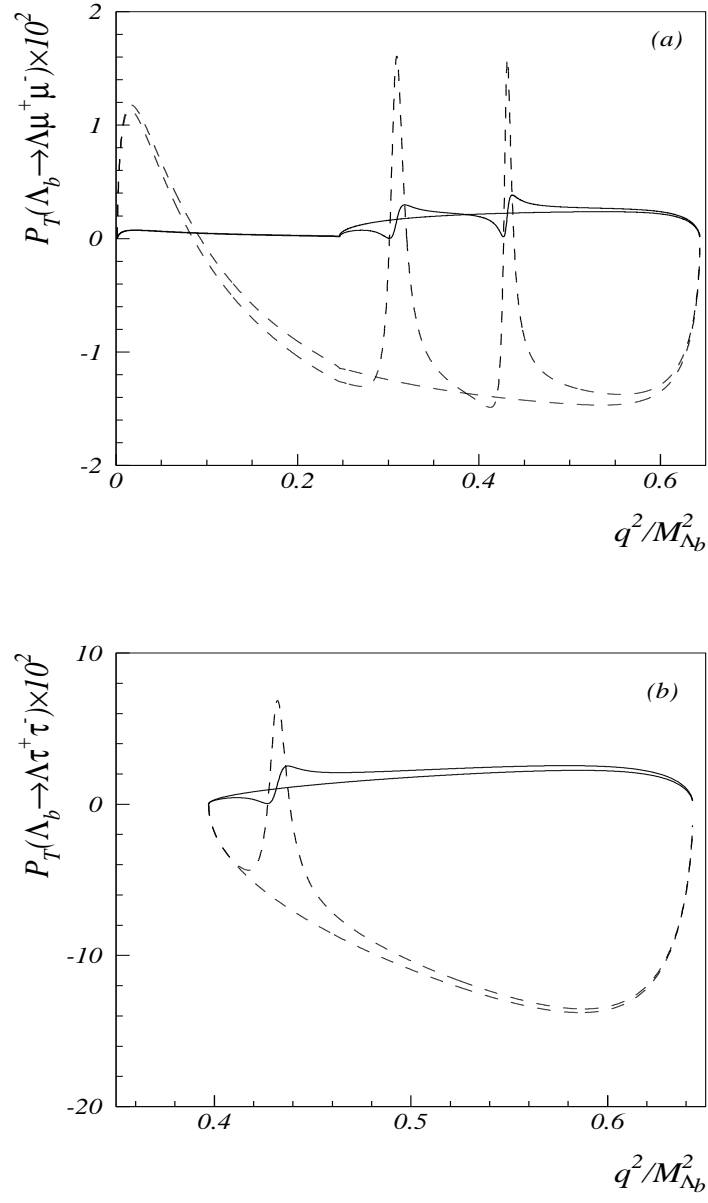


Figure 7: Same as Figure 5 but for the transverse polarization asymmetries.

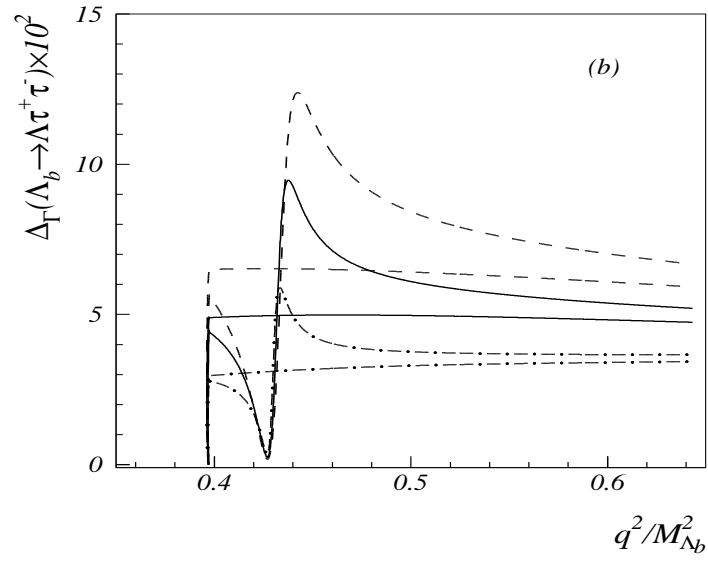
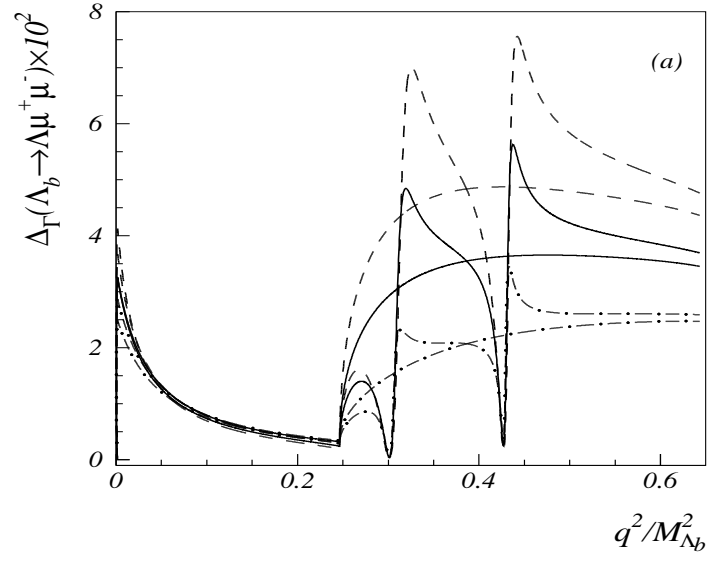


Figure 8: Same as Figure 5 but for Δ_Γ .

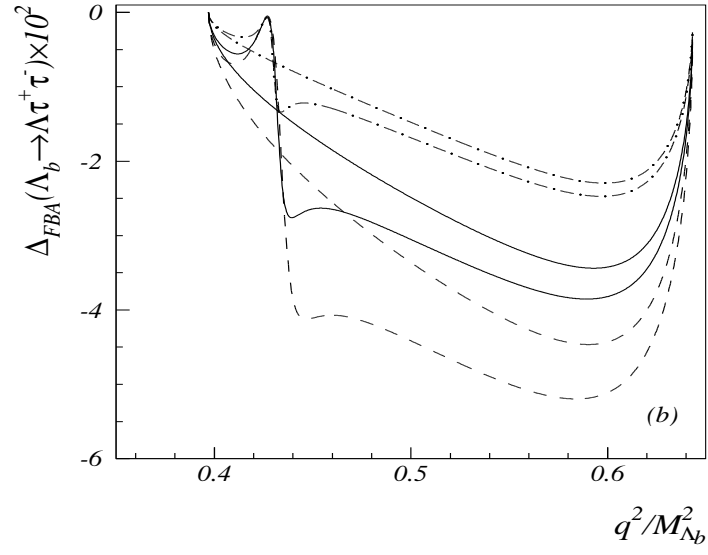
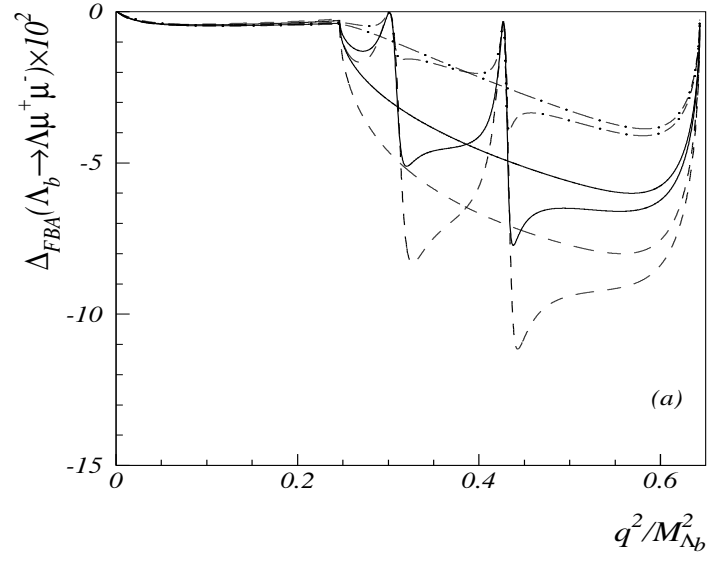


Figure 9: Same as Figure 5 but for Δ_{FBA} .

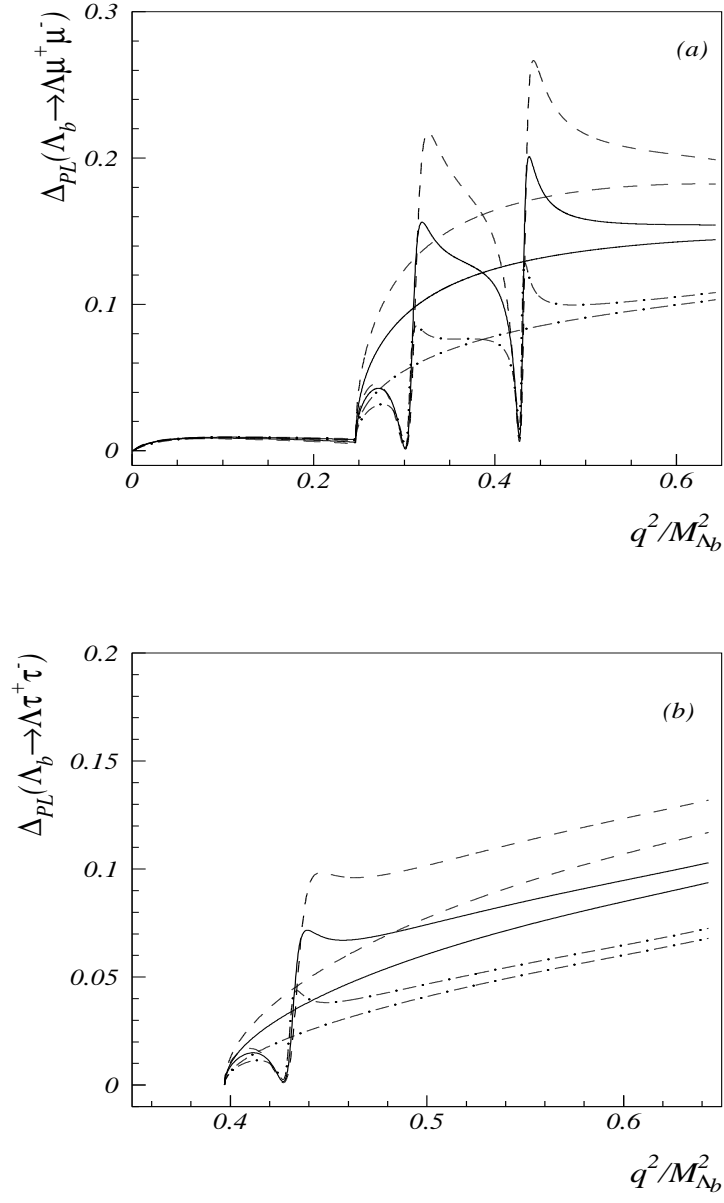


Figure 10: Same as Figure 5 but for Δ_{PL} .

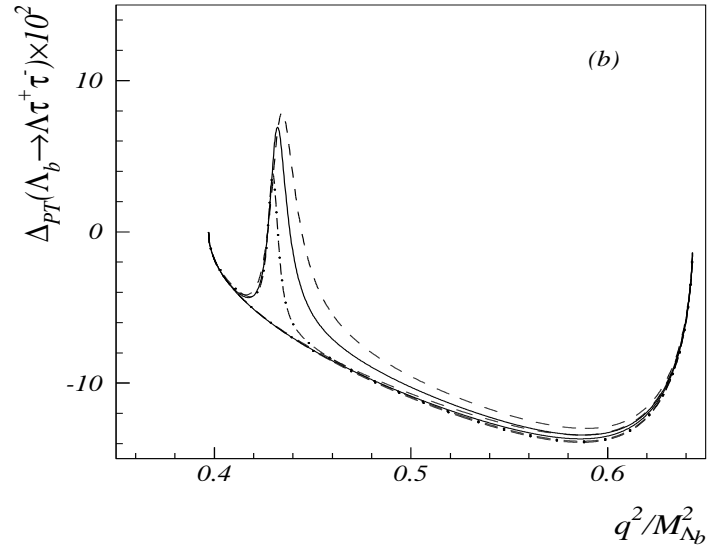
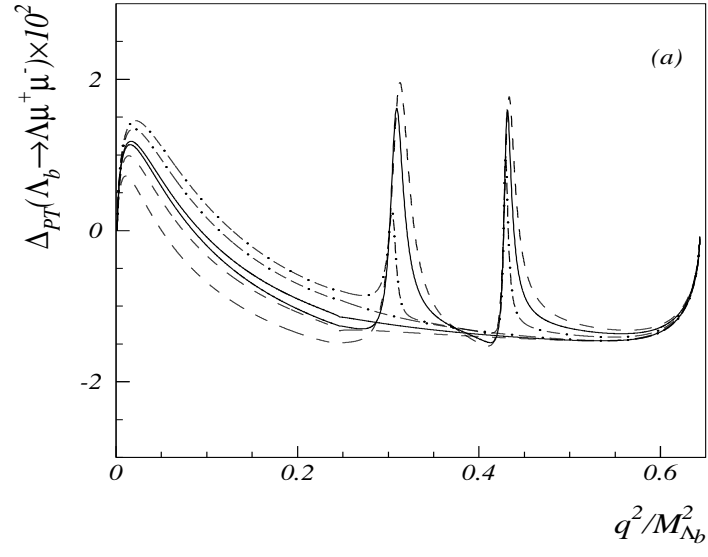


Figure 11: Same as Figure 5 but for Δ_{P_T} .